TUM/T39-99-15
hep-ph/9907373

Gluons and the η'–nucleon coupling constant *

Steven D. Bass

Physik Department, Technische Universität München,
D-85747 Garching, Germany

Abstract

We derive the effective chiral Lagrangian for low-energy η–nucleon and η'–nucleon interactions and show that gluonic degrees of freedom, via the axial $U_A(1)$ anomaly, induce a contact term in the $pp \rightarrow pp\eta$ and $pp \rightarrow pp\eta'$ reactions. We then discuss the consequences for the extraction of $g_{\eta' NN}$ from experimental data.

*Work supported in part by BMBF.

1 Introduction

Polarised deep inelastic scattering [1, 2] and η' physics [3] provide complementary windows on dynamics induced by the axial $U_A(1)$ anomaly [4] in QCD. The flavour-singlet Goldberger-Treiman relation [5] relates the flavour-singlet axial-charge $g_A^{(0)}$ measured in polarised deep inelastic scattering to the η'–nucleon coupling constant $g_{\eta' NN}$. The large mass of the η' and the small value of $g_A^{(0)}$ extracted from deep inelastic scattering point to substantial violations of the OZI rule [6] in the flavour-singlet $J^P = 1^+$ channel [7]. In this paper we discuss the role of gluonic degrees of freedom in the low-energy η'–nucleon interaction. We derive the effective chiral Lagrangian for η'–nucleon interactions and explain why gluons induce a contact term in the low-energy $pp \to pp\eta'$ reaction. The strength of this contact term is, in part, related to the amount of spin carried by polarised gluons in a polarised proton. The total cross section for $pp \to pp\eta'$ near threshold has been measured at COSY [8] and SATURNE [9]. New measurements could follow from an upgraded CELSIUS machine. The η' photoproduction process is being studied in experiments at ELSA [10] and Jefferson Laboratory [11].

We begin in Section 2 with a brief review of the effective Lagrangian [12, 13, 14] for low-energy η'–meson interactions. In Section 3 we extend this theory to include η'–nucleon coupling. Finally, in Section 4, we discuss the possible size of gluonic effects in the η'–nucleon interaction and the consequences for the extraction of $g_{\eta' NN}$ from experimental data.

2 The low-energy effective Lagrangian

Starting in the meson sector, the effective Lagrangian [12, 13, 14] for low-energy QCD is constructed as follows. We begin with the SU(3) chiral Lagrangian

$$\mathcal{L} = \frac{F_\pi^2}{4} \mathrm{Tr}(\partial^\mu U \partial_\mu U^\dagger) + \frac{F_\pi^2}{4} \mathrm{Tr} M \left(U + U^\dagger \right) \tag{1}$$

where U is the unitary meson matrix and $M = \mathrm{diag}[m_\pi^2, m_\pi^2, (2m_K^2 - m_\pi^2)]$ is the meson mass matrix — for a review see [15]. The matrix U is extended to include a flavour-singlet Goldstone boson [16]:

$$U = \exp \left(i \frac{\phi}{F_\pi} + i \sqrt{\frac{2}{3}} \frac{\eta_0}{F_0} \right). \tag{2}$$

Here

$$\phi = \sqrt{2} \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta_8 & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta_8 & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta_8 \end{pmatrix} \tag{3}$$

denotes the octet of would-be Goldstone bosons associated with spontaneous chiral $SU(3)_L \otimes SU(3)_R$ breaking and η_0 is the singlet boson. The pion decay constant

$F_\pi = 92.4$MeV; F_0 renormalises the flavour-singlet decay constant — see Eq.(8) below.

Next, one introduces $U_A(1)$ terms to generate a gluonic mass term for the η_0 and to reproduce the anomaly [17, 18] in the divergence of the gauge-invariantly renormalised flavour-singlet axial-vector current

$$J_{\mu 5} = \left[\bar{u}\gamma_\mu\gamma_5 u + \bar{d}\gamma_\mu\gamma_5 d + \bar{s}\gamma_\mu\gamma_5 s\right]^{\mu^2}_{GI}, \tag{4}$$

viz.

$$\partial^\mu J_{\mu 5} = \sum_{k=1}^{f} 2i\left[m_k \bar{q}_k\gamma_5 q_k\right]_{GI} + N_f\left[\frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G}^{\mu\nu}\right]^{\mu^2}_{GI}. \tag{5}$$

Here $N_f = 3$ is the number of light flavours, the subscript GI denotes gauge invariant renormalisation and the superscript μ^2 denotes the renormalisation scale. The Adler-Bardeen theorem [19] states that the anomaly on the right hand side of Eq.(5) is not renormalised to all orders in perturbation theory. We use this theorem to constrain the possible $U_A(1)$ breaking terms in the effective Lagrangian.

The low-energy effective Lagrangian [13, 14] is

$$\begin{aligned}\mathcal{L}_{\rm m} = \quad & \frac{F_\pi^2}{4}{\rm Tr}(\partial^\mu U\partial_\mu U^\dagger) + \frac{F_\pi^2}{4}{\rm Tr}M\left(U + U^\dagger\right) \\ + \quad & b\frac{F_\pi^2}{4}\left({\rm Tr}U^\dagger\partial_\mu U\right)^2 + \frac{1}{2}iQ{\rm Tr}\left[\log U - \log U^\dagger\right] + \frac{1}{aF_0^2}Q^2.\end{aligned} \tag{6}$$

The gluonic term $Q = \frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G}^{\mu\nu}$ is treated as a background field. It has no kinetic term [1] and mixes with the η_0 to generate a gluonic mass term for the η' — see below. The term proportional to b is fourth order in U but gives a second order term in η_0. For simplicity we omit fourth order terms in the meson fields. Their inclusion is straightforward but of no primary influence in our context. The most general low-energy effective Lagrangian involves a $U_A(1)$ invariant polynomial in Q^2. Higher-order terms in Q^2 become important when we consider scattering processes involving more than one η' [20].

The $U_A(1)$ transformation for the effective Lagrangian (6) is defined [3] by $U \to \exp(2i\beta)\, U$; Q is treated as $U_A(1)$ invariant. The flavour-singlet axial-vector current

$$J_{\mu 5} = \sqrt{6}F_{\rm singlet}\partial_\mu\eta_0 \tag{7}$$

with

$$F_{\rm singlet} = \frac{F_\pi^2}{F_0}\left(1 - 3b\right) \tag{8}$$

satisfies the anomalous divergence equation $\partial^\mu J_{\mu 5} = N_f Q +$ mass terms. The flavour-singlet decay constant is renormalised relative to F_π by gluonic intermediate states ($q\bar{q} \to gg \to q\bar{q}$). The parameters F_0 and b describe two distinct renormalisation effects: F_0 contributes to $F_{\rm singlet}$ and also the meson masses (through its appearance in ${\rm Tr}M(U + U^\dagger)$) whereas b contributes only to $F_{\rm singlet}$. Assuming a continuous large N_c limit,

$$\lim_{N_c\to\infty} F_{\rm singlet} = F_\pi \tag{9}$$

[1] In QCD $Q = \partial^\mu K_\mu$ where K_μ is the gauge-dependent gluonic Chern-Simons current. The reader may wish to think of Q^2 as a kinetic term for K_μ.

in the chiral limit.

The gluonic term Q can be eliminated from the effective Lagrangian through its equation of motion. Expanding to $\mathcal{O}(p^2)$ in momentum and keeping finite quark masses one finds

$$\begin{aligned}
\mathcal{L}_{\rm m} = \quad & \frac{1}{2}\partial^\mu \pi_a \partial_\mu \pi_a + \frac{1}{2}\partial_\mu \eta_0 \partial^\mu \eta_0 \left(\frac{F_\pi}{F_0}\right)^2 (1-3b) \\
- \quad & \frac{3a}{2}\eta_0^2 \\
- \quad & \frac{1}{2}m_\pi^2\Big(2\pi^+\pi^- + \pi_0^2\Big) - m_K^2\Big(K^+K^- + K^0\bar{K}^0\Big) - \frac{1}{2}\Big(\frac{4}{3}m_K^2 - \frac{1}{3}m_\pi^2\Big)\eta_8^2 \\
- \quad & \frac{1}{2}\Big(\frac{2}{3}m_K^2 + \frac{1}{3}m_\pi^2\Big)\left(\frac{F_\pi}{F_0}\right)^2 \eta_0^2 + \frac{4}{3\sqrt{2}}\Big(m_K^2 - m_\pi^2\Big)\left(\frac{F_\pi}{F_0}\right)\eta_8\eta_0
\end{aligned} \tag{10}$$

where ϕ_a with $a = 1, ..., 8$ refers to the octet Goldstone boson fields. In the chiral limit $M = 0$ gluons contribute a finite mass

$$m_{\eta_0}^2 = \frac{3a}{\left(\frac{F_\pi}{F_0}\right)^2 (1-3b)} \tag{11}$$

to the singlet η_0. The masses of the physical η and η' mesons are found by diagonalising the (η_8, η_0) mass matrix which follows from Eq.(10).

Henceforth, we take $F_0 \sim 0.1$GeV [21] and $\tilde{m}_{\eta_0}^2 \equiv 3a \sim 1\text{GeV}^2$.

It is worthwhile to comment on the behaviour of $\mathcal{L}_{\rm m}$ when we take N_c, the number of colours, to infinity. In QCD the axial anomaly decouples as $1/N_c$ when we take $N_c \to \infty$. This is reflected in the equation for the η_0 mass squared if we take $a \propto 1/N_c$ [12]. Phenomenologically, the large mass of the η' ($m_{\eta'} \sim 1$GeV) means that OZI and large N_c are not always good approximations in the $U_A(1)$ channel. A second source of OZI violation is the non-vanishing anomalous dimension [22, 23, 24] of the flavour-singlet axial vector current (4). In QCD $J_{\mu5}$ satisfies the renormalisation group equation

$$J_{\mu5}(\lambda) = J_{\mu5}(\infty)/E[\alpha_s(\lambda)] \tag{12}$$

where

$$E[\alpha_s(\lambda)] = \exp \int_0^{\alpha_s(\lambda)} d\tilde{\alpha}_s\, \gamma(\tilde{\alpha}_s)/\beta(\tilde{\alpha}_s). \tag{13}$$

Here $\gamma(\alpha_s)$ $(= f\frac{\alpha_s^2}{\pi^2} + \mathcal{O}(\alpha_s^3))$ is the (two loop) non-zero anomalous dimension of $J_{\mu5}$ and $\beta(\alpha_s)$ is the QCD beta function. We are free to choose the QCD coupling $\alpha_s(\lambda)$ at either a hard or a soft scale λ. If we work to $O(\alpha_s^2)$ in perturbation theory, then $E[\alpha_s] \sim 0.84$ when $\alpha_s \sim 0.6$, typical of the infrared – see eg. [1]. At least in perturbation theory, $E[\alpha_s]$ remains close to the OZI value $E[\alpha_s] = 1$ — in contrast to $m_{\eta_0}^2$ which exhibits large OZI violation. For the meson theory (6) the renormalisation group factor $E[\alpha_s]$ can be absorbed in the parameter b.

3 The η'–nucleon interaction

The low-energy effective Lagrangian (6) is readily extended to include η'–nucleon coupling. For simplicity we work in the chiral limit. The chiral SU(3) meson-baryon

coupling Lagrangian is

$$\begin{aligned}\mathcal{L}_{\rm mB} = & \quad {\rm Tr}\overline{B}(i\gamma_\mu D^\mu - m)B \\ & + \quad F\ {\rm Tr}\Big(\overline{B}\gamma_\mu\gamma_5[a^\mu, B]_-\Big) + D\ {\rm Tr}\Big(\overline{B}\gamma_\mu\gamma_5\{a^\mu, B\}_+\Big)\end{aligned} \tag{14}$$

where

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda & n \\ \Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}}\Lambda \end{pmatrix} \tag{15}$$

denotes the baryon octet and m denotes the baryon mass. In Eq.(14) $D_\mu = \partial_\mu - i v_\mu$ is the chiral covariant derivative, $v_\mu = -\frac{i}{2}\Big(\xi^\dagger\partial_\mu\xi + \xi\partial_\mu\xi^\dagger\Big)$ and $a_\mu = -\frac{i}{2}\Big(\xi^\dagger\partial_\mu\xi - \xi\partial_\mu\xi^\dagger\Big)$ where $\xi = U^{\frac{1}{2}}$. The SU(3) couplings are $F = 0.459 \pm 0.008$ and $D = 0.798 \pm 0.008$ [25]. The Pauli-principle forbids any flavour-singlet $J^P = \frac{1}{2}^+$ ground-state baryon degenerate with the baryon octet B.

We work to $O(Q^2)$ in the gluonic field and add the leading $U_A(1)$ invariant terms

$$\begin{aligned}\mathcal{L}_{\rm U_A(1)\ mB} = & \quad \frac{i}{3}K\ {\rm Tr}\Big(\overline{B}\gamma_\mu\gamma_5 B\Big){\rm Tr}\Big(U^\dagger\partial^\mu U\Big) \\ & - \quad \frac{\mathcal{G}_{QNN}}{2m}\partial^\mu Q {\rm Tr}\Big(\overline{B}\gamma_\mu\gamma_5 B\Big) + \frac{\mathcal{C}}{F_0^4}Q^2{\rm Tr}\Big(\overline{B}B\Big).\end{aligned} \tag{16}$$

Here, we have defined $\mathcal{G}_{QNN}$ and $\mathcal{C}$ so that they both have mass dimension -3. Some comment is warranted about the momentum scale involved in the coefficient multiplying the $Q^2\ {\rm Tr}(\bar{B}B)$ term. There are two natural choices. The form $\mathcal{C}/F_0^4$ written in Eq.(16) is motivated by the fact that $F_0 \sim F_\pi$ is relatively free of OZI violation. Alternatively, motivated by the coefficient of Q^2 in Eq.(6), one might choose to replace $\mathcal{C}/F_0^4$ in Eq.(16) by $\mathcal{C}'/aF_0^2$ with $\mathcal{C}' = \mathcal{C}a/F_0^2$. In this case the OZI violation partially cancels in the ratio $\mathcal{C}'/a = 3\mathcal{C}'/\tilde{m}^2_{\eta_0}$. Note the relative size of the two OZI parameters $\mathcal{C}' \sim 33\ \mathcal{C}$. We proceed keeping the $\mathcal{C}/F_0^4$ coefficient in Eq.(16) and return to this point in Section 4.

Extra $U_A(1)$ breaking terms of the form ${\rm Tr}\Big(\log U - \log U^\dagger\Big)\ Q\ {\rm Tr}\Big(\overline{B}B\Big)$ or $\Big({\rm Tr}(\log U - \log U^\dagger)\Big)^2{\rm Tr}\Big(\overline{B}B\Big)$ are excluded by the Adler-Bardeen theorem. Such terms would correspond to extra $U_A(1)$ breaking in the divergence equation for $J_{\mu 5}$ — in contradiction to the non-renormalisation of the axial anomaly in Eq.(5). The divergence of the flavour-singlet axial-vector current in the effective theory consists of the mass terms we obtain with $Q = \frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G}^{\mu\nu}$ turned off plus a single anomaly factor of $N_f Q$.

Since we are interested in the $pp \to pp\eta'$ reaction we also include possible fourth order (contact) terms in the baryon fields:

$$\mathcal{L}_{\rm m2B} = \quad \Big[\lambda_F F\ {\rm Tr}\Big(\overline{B}\gamma_\mu\gamma_5[a^\mu, B]_-\Big) + \lambda_D D\ {\rm Tr}\Big(\overline{B}\gamma_\mu\gamma_5\{a^\mu, B\}_+\Big) \tag{17}$$

$$+\frac{i}{3}\lambda_K K\ \mathrm{Tr}\Big(\overline{B}\gamma_\mu\gamma_5 B\Big)\mathrm{Tr}\Big(U^\dagger\partial^\mu U\Big)$$
$$+\lambda_Q\frac{\mathcal{G}_{QNN}}{2m}\ Q\ \partial^\mu\mathrm{Tr}\Big(\overline{B}\gamma_\mu\gamma_5 B\Big)\Big]\ \mathrm{Tr}\Big(\bar{B}B\Big)$$

where λ_F, λ_D, λ_K and λ_Q are new parameters; $\mathcal{L}_{\mathrm{m2B}}$ is $U_A(1)$ invariant for the same reason as $\mathcal{L}_{\mathrm{U_A(1)\ mB}}$. Putting things together, our effective Lagrangian is

$$\mathcal{L}=\mathcal{L}_{\mathrm{m}}+\mathcal{L}_{\mathrm{mB}}+\mathcal{L}_{\mathrm{U_A(1)\ mB}}+\mathcal{L}_{\mathrm{m2B}}. \tag{18}$$

Let us consider the Q dependent terms in more detail. The Q dependent part of the effective Lagrangian (18) is

$$\begin{aligned}\mathcal{L}_Q\ &=\ \frac{1}{2}i\ Q\ \mathrm{Tr}\Big[\log U-\log U^\dagger\Big]+\frac{1}{aF_0^2}\ Q^2 \qquad (19)\\ &+\ \frac{\mathcal{G}_{QNN}}{2m}\ Q\ \partial^\mu\mathrm{Tr}\Big(\bar{B}\gamma_\mu\gamma_5 B\Big)\ \Big\{1+\lambda_Q\mathrm{Tr}\bar{B}B\Big\}+\frac{\mathcal{C}}{F_0^4}\ Q^2\ \mathrm{Tr}\Big(\bar{B}B\Big),\end{aligned}$$

which yields the following equation of motion for Q:

$$\begin{aligned}&\frac{2}{aF_0^2}\ \Big(1+\frac{\mathcal{C}}{F_0^2}\ a\ \mathrm{Tr}(\overline{B}B)\Big)\ Q \qquad (20)\\ &\quad=-\Big(\frac{1}{2}i\mathrm{Tr}\Big[\log U-\log U^\dagger\Big]+\frac{\mathcal{G}_{QNN}}{2m}\partial^\mu\mathrm{Tr}\Big(\overline{B}\gamma_\mu\gamma_5 B\Big)\Big\{1+\lambda_Q\mathrm{Tr}\bar{B}B\Big\}\Big).\end{aligned}$$

We substitute for Q in $\mathcal{L}_Q$ to obtain

$$\begin{aligned}\mathcal{L}_Q\ =\ -\frac{1}{12}\tilde{m}_{\eta_0}^2\ \Big[\ &6\eta_0^2\ +\ \frac{\sqrt{6}}{m}\ \mathcal{G}_{QNN}\ F_0\ \partial^\mu\eta_0\ \mathrm{Tr}\Big(\bar{B}\gamma_\mu\gamma_5 B\Big) \qquad (21)\\ &-\ \mathcal{G}_{QNN}^2\ F_0^2\ \Big(\mathrm{Tr}\bar{B}\gamma_5 B\Big)^2\ -\ 2\ \mathcal{C}\ \frac{\tilde{m}_{\eta_0}^2}{F_0^2}\ \eta_0^2\ \mathrm{Tr}\Big(\bar{B}B\Big)\\ &+\ \frac{\sqrt{6}}{mF_0}\ \mathcal{G}_{QNN}\ \Big\{\frac{1}{3}\mathcal{C}\tilde{m}_{\eta_0}^2-F_0^2\lambda_Q\Big\}\ \eta_0\ \partial^\mu\mathrm{Tr}\Big(\bar{B}\gamma_\mu\gamma_5 B\Big)\ \mathrm{Tr}\Big(\bar{B}B\Big)+...\Big]\end{aligned}$$

where we have replaced a by $\frac{1}{3}\tilde{m}_{\eta_0}^2$.

The Lagrangian (21) has three contact terms associated with the gluonic potential in Q. First, we observe the contact term

$$\mathcal{L}^{(1)}_{\mathrm{contact}}=\frac{1}{12}\ \tilde{m}_{\eta_0}^2\ \mathcal{G}_{QNN}^2\ F_0^2\ \Big(\mathrm{Tr}\bar{B}\gamma_5 B\Big)^2 \tag{22}$$

in the nucleon-nucleon interaction which was discovered by Schechter et al [26]. Second, we find two new contact terms associated with the axial anomaly:

$$\mathcal{L}^{(2)}_{\mathrm{contact}}=-\frac{\sqrt{6}}{12mF_0}\mathcal{G}_{QNN}\tilde{m}_{\eta_0}^2\ \Big\{\frac{1}{3}\mathcal{C}\tilde{m}_{\eta_0}^2-F_0^2\lambda_Q\Big\}\ \eta_0\ \partial^\mu\mathrm{Tr}\Big(\bar{B}\gamma_\mu\gamma_5 B\Big)\ \mathrm{Tr}\Big(\bar{B}B\Big) \tag{23}$$

and

$$\mathcal{L}^{(3)}_{\mathrm{contact}}=\frac{1}{6F_0^2}\ \mathcal{C}\ \tilde{m}_{\eta_0}^4\ \eta_0^2\ \mathrm{Tr}\Big(\bar{B}B\Big). \tag{24}$$

The contact interaction $\mathcal{L}^{(2)}_{\rm contact}$ contributes to the low-energy $pp \to pp\eta'$ reaction; $\mathcal{L}^{(3)}_{\rm contact}$ is relevant to $\eta' N \to \eta' N$. The three terms (22-24) describe short distance interactions. The term $\mathcal{L}^{(2)}_{\rm contact}$ yields an additional contribution to the cross-section for $pp \to pp\eta'$ reaction which is extra to the long distance contributions associated with meson exchange models [27, 28]. The contact terms $\mathcal{L}^{(j)}_{\rm contact}$ are proportional to $\tilde{m}^2_{\eta_0}$ $(j=1,2)$ and $\tilde{m}^4_{\eta_0}$ $(j=3)$ which vanish in the formal OZI limit. Phenomenologically, the large masses of the η and η' mesons means that there is no reason, a priori, to expect the $\mathcal{L}^{(j)}_{\rm contact}$ to be small.

The total contact term in the $pp \to pp\eta_0$ reaction is obtained by combining $\mathcal{L}^{(2)}_{\rm contact}$ with the contact terms coming from the Q independent part of Eq.(17):

$$\begin{aligned}\mathcal{L}^{(4)}_{\rm contact} = \Bigg[\ &- \ \sqrt{\frac{2}{3}}\Big\{ \frac{\lambda_D D}{F_0} + \frac{\lambda_K K}{F_0} \Big\} \qquad\qquad (25) \\ &- \ \frac{\sqrt{6}}{12 m F_0} \mathcal{G}_{QNN} \tilde{m}^2_{\eta_0} \ \Big\{ \frac{1}{3} \mathcal{C} \tilde{m}^2_{\eta_0} - F_0^2 \lambda_Q \Big\} \Bigg] \eta_0 \ \partial^\mu {\rm Tr}\Big(\bar{B} \gamma_\mu \gamma_5 B \Big) \ {\rm Tr}\Big(\bar{B} B \Big).\end{aligned}$$

If we substitute $\mathcal{C} = \mathcal{C}' a / F_0^2$ then the factor $\{\frac{1}{3}\mathcal{C}\tilde{m}^2_{\eta_0} - F_0^2 \lambda_Q\}$ becomes $F_0^2\{\mathcal{C}' - \lambda_Q\}$. We now discuss plausible values for K, $\mathcal{G}_{QNN}$, $\mathcal{C}$ and the strength of the contact term in Eq.(25).

4 η' nucleon coupling constants: how many ?

The physical η'–nucleon coupling constant is read off from the coefficient of $\partial^\mu \eta_0 {\rm Tr} \bar{B} B$ in the effective Lagrangian *after* we have eliminated Q. One finds

$$g_{\eta_0 NN} = \sqrt{\frac{2}{3}} \frac{m}{F_0} \Big(2D + 2K + \mathcal{G}_{QNN} F_0^2 \frac{\tilde{m}^2_{\eta_0}}{2m} \Big). \qquad (26)$$

This $g_{\eta_0 NN}$ is the η_0–nucleon coupling constant which will enter the extension of the coupled channels analysis [29] to η' photoproduction.

The flavour-singlet Goldberger-Treiman relation for QCD was derived by Shore and Veneziano [5]. In the notation of Eq.(16) one finds

$$m g_A^{(0)} = \sqrt{\frac{3}{2}} F_0 \Big(g_{\eta_0 NN} - g_{QNN} \Big) \qquad (27)$$

where $g_{QNN} = \sqrt{\frac{1}{6}} \mathcal{G}_{QNN} F_0 \tilde{m}^2_{\eta_0}$. If gluonic degrees of freedom were turned off then the right hand side of Eq.(27) would be equal to $F_0 g_{\eta_0 NN}$. The value of $g_A^{(0)}$ extracted from polarised deep inelastic scattering is [1, 2]

$$g_A^{(0)}\Big|_{\rm pDIS} = 0.2 - 0.35. \qquad (28)$$

The OZI prediction $g_A^{(0)} \simeq 0.6$ [30] would follow if polarised strange quarks and gluons were not important in the nucleon's internal spin structure. *If* we attribute the difference between $g_A^{(0)}|_{\rm pDIS}$ and the OZI value 0.6 to the gluonic correction $-\sqrt{\frac{3}{2}} F_0 g_{QNN}$ in Eq.(27), *then* we find $g_{QNN} \sim 2.45$ and $g_{\eta_0 NN} \sim 4.9$ with $F_0 \sim$

0.1GeV. These values correspond to $K \sim -0.65$ and $\mathcal{G}_{QNN} \sim +60\text{GeV}^{-3}$ if we take $\tilde{m}^2_{\eta_0} \sim 1\text{GeV}^2$ and substitute into Eq.(26). The coupling constant g_{QNN} is, in part, related [5] to the amount of spin carried by polarised gluons in a polarised proton.

How important is the contact interaction $\mathcal{L}^{(4)}_{\text{contact}}$ in the $pp \rightarrow pp\eta'$ reaction ?

The T-matrix for η' production in proton-proton collisions, $p_1(\vec{p}) + p_2(-\vec{p}) \rightarrow p + p + \eta'$, at threshold in the centre of mass frame is

$$\mathrm{T}^{\mathrm{cm}}_{\mathrm{th}}(pp \rightarrow pp\eta') = \mathcal{A}\Big[i(\vec{\sigma}_1 - \vec{\sigma}_2) + \vec{\sigma}_1 \mathrm{x} \vec{\sigma}_2\Big].\vec{p} \tag{29}$$

where $\mathcal{A}$ is the (complex) threshold amplitude for η' production. Measurements of the total cross-section for $pp \rightarrow pp\eta'$ have been published by COSY [8] and SATURNE [9] at centre of mass energies 1.5, 1.7, 2.9 and 3.7 MeV above threshold (COSY) and 4.1 and 8.3 MeV above threshold (SATURNE). If taken at face value, the four COSY data points are best described by pure three-body phase space. If one ignores the COSY data at the two lowest energies, then the remaining four data points are well described using the model of Bernard et al. [31] treating the pp final state interaction in effective range approximation. In this model, which also provides a good description of $pp \rightarrow pp\pi^0$ and $pp \rightarrow pp\eta$ close to threshold, one finds a best fit to the measured total cross-section data with

$$|\mathcal{A}| = 0.21 \text{ fm}^4 \tag{30}$$

with $\chi^2 = 2.4$. The present (total cross-section only) data on $pp \rightarrow pp\eta'$ is insufficient to distinguish between possible production mechanisms involving the contact term (25) and meson exchange models. To estimate how strong the contact term must be in order to make an important contribution to the measured $pp \rightarrow pp\eta'$ cross-section, let us consider the extreme scenario where the value of $|\mathcal{A}|$ in Eq.(30) is saturated by the contact term (25).

For the values $K \simeq -0.65$ and $\mathcal{G}_{QNN} \simeq 60\text{GeV}^{-3}$ — see below Eq.(28) — the contact term (25) becomes

$$\mathcal{L}^{(4)}_{\text{contact}} = \Big(13.1\lambda_D - 10.6\lambda_K - 81.6\mathcal{C} + 2.4\lambda_Q\Big)\ \eta_0\ \mathrm{Tr}\Big(\bar{B}i\gamma_5 B\Big)\ \mathrm{Tr}\Big(\bar{B}B\Big). \tag{31}$$

The success of meson exchange models in describing the low-energy $pp \rightarrow p\Lambda K$ and $pp \rightarrow pp\eta$ reactions [28, 9] suggests that it may be reasonable to take the coefficients λ_D and λ_K of the OZI preserving processes small. The large coefficient of $\mathcal{C}$ in Eq.(31) is induced by the OZI violation in $\tilde{m}^2_{\eta_0}$ and $\mathcal{G}_{QNN}$ (assuming that g_{QNN} is indeed responsible for the small value of $g_A^{(0)}|_{\text{pDIS}}$ extracted from polarised deep inelastic scattering). If we saturate $|\mathcal{A}|$ in Eq.(30) by the OZI violating term proportional to $\mathcal{C}$ then we obtain $\mathcal{C} \sim 1.8\text{GeV}^{-3}$. Alternatively, if we use $\tilde{m}^2_{\eta_0}$ to set the scale for $Q^2\mathrm{Tr}(\bar{B}B)$ and take $\mathcal{C}/F_0^4 \mapsto \mathcal{C}'/aF_0^2$ in Eqs.(16) and (25), then $\mathcal{C}'$ and λ_Q enter Eq.(31) with the same coefficient 2.4. Saturating $|\mathcal{A}|$ in Eq.(30) with the gluonic contact term proportional to $\{\mathcal{C}' - \lambda_Q\}$ yields $\{\mathcal{C}' - \lambda_Q\} \sim 60\text{GeV}^{-3}$. The OZI violating parameters $\mathcal{C} \sim 1.8\text{GeV}^{-3}$ and $\{\mathcal{C}' - \lambda_Q\} \sim 60\text{GeV}^{-3}$ compare with $\mathcal{G}_{QNN} \sim 60\text{GeV}^{-3}$.

In QCD the large strange quark mass generates substantial η–η' mixing. Working in the one-mixing-angle scheme [21] one finds

$$|\eta_0\rangle = \cos\theta\ |\eta'\rangle - \sin\theta\ |\eta\rangle \tag{32}$$

with $\theta \simeq -20$ degrees. The ratio of the moduli of the contact amplitudes (25) for $pp \to pp\eta'$ and $pp \to pp\eta$ is proportional to $\cos\theta : \sin\theta = 0.94 : 0.34$. The contact interaction (25) is likely to play a much more prominant role in the cross-section for $pp \to pp\eta'$ than $pp \to pp\eta$. In their analysis of the SATURNE data on $pp \to pp\eta'$ Hibou et al. [9] found that a one-pion exchange model adjusted to fit the S-wave contribution to the $pp \to pp\eta$ cross-section near threshold yields predictions about 30% below the measured $pp \to pp\eta'$ total cross-section. The gluonic contact term (23) is a candidate for additional, potentially important, short range interaction.

Gluonic $U_A(1)$ degrees of freedom induce several "η'–nucleon coupling constants": $g_{\eta_0 NN}$, $\mathcal{G}_{QNN}$ and $\mathcal{C}$. Different combinations of these coupling constants are relevant to different η' production processes and to the flavour-singlet Goldberger-Treiman relation. Testing the sensitivity of η'–nucleon interactions to the gluonic terms in the effective chiral Lagrangian for low-energy QCD will teach us about the role of gluons in chiral dynamics.

Acknowledgements

I thank B. Borasoy, G. Fäldt, N. Kaiser and W. Weise for helpful discussions.

References

[1] S. D. Bass, hep-ph/9902280, Eur. Phys. J **A5** (1999) 17.

[2] R. Windmolders, hep-ph/9905505.

[3] G. Christos, Phys. Rept. **116** (1984) 251.

[4] G.M. Shore, Zuoz lecture, hep-ph/9812354.

[5] G. Veneziano, Mod. Phys. Lett. **A4** (1989) 1605;
G.M. Shore and G. Veneziano, Nucl. Phys. **B381** (1992) 23.

[6] S. Okubo, Phys. Lett. **5** (1963) 165;
G. Zweig, CERN report No. 8419/TH412 (1964);
J. Iizuka, Prog. Theor. Phys. Suppl. **37-38** (1966) 21.

[7] G. Veneziano, Okubofest lecture, CERN preprint TH-5840 (1990).

[8] The COSY-11 Collaboration (P. Moskal et al.), Phys. Rev. Lett. **80** (1998) 3202.

[9] SATURNE, F. Hibou et al., Phys. Lett. **B438** (1998) 41.

[10] The SAPHIR Collaboration, R. Plötzke et al., Phys. Lett. **B444** (1998) 555.

[11] CEBAF experiments E-89-039 (S. Dytman et al.) and E-91-008 (B.G. Ritchie et al.).

[12] E. Witten, Nucl. Phys. **B156** (1979) 269;
G. Veneziano, Nucl. Phys. **B159** (1979) 213.

[13] P. Di Vecchia and G. Veneziano, Nucl. Phys. **B171** (1980) 253.

[14] C. Rosenzweig, J. Schechter and C.G. Trahern, Phys. Rev. **D21** (1980) 3388;
P. Nath and R. Arnowitt, Phys. Rev. **D23** (1981) 473.

[15] W. Weise, Les Houches lectures, in *Trends in Nuclear Physics, 100 Years Later*, (Elsevier Science, 1998)

[16] S. Weinberg, *The Quantum Theory of Fields*, Vol. II (Cambridge UP, 1996).

[17] S.L. Adler, Phys. Rev. **177** (1969) 2426.

[18] J.S. Bell and R. Jackiw, Nuovo Cimento **60A** (1969) 47.

[19] S.L. Adler and W.A. Bardeen, Phys. Rev. **182** (1969) 1517.

[20] P. Di Vecchia, F. Nicodemi, R. Pettorino and G. Veneziano, Nucl. Phys. **B181** (1981) 318.

[21] F.J. Gilman and R. Kauffman, Phys. Rev. **D36** (1987) 2761; (E) **D37** (1988) 3348.

[22] R. Köberle and N.K. Nielsen, Phys. Rev. **D8** (1973) 660;
R.J. Crewther, S.-S. Shei and T.-M. Yan, Phys. Rev. **D8** (1973) 3396;
R.J. Crewther and N.K. Nielsen, Nucl. Phys. **B87** (1975) 52.

[23] R.J. Crewther, Schlading lectures, Acta Physica Austriaca Suppl. **19** (1978) 47, section 7.

[24] J. Kodaira, Nucl. Phys. **B165** (1980) 129.

[25] F.E. Close and R.G. Roberts, Phys. Lett. **B316** (1993) 165.

[26] J. Schechter, V. Soni, A. Subbaraman and H. Weigel, Mod. Phys. Lett. **A5** (1990) 2543.

[27] R. Machleidt, K. Holinde and Ch. Elster, Phys. Rept. **149** (1987) 1.

[28] G. Fäldt and C. Wilkin, Z Physik **A357** (1997) 241.

[29] N. Kaiser, T. Waas and W. Weise, Nucl. Phys. **A612** (1997) 297.

[30] J. Ellis and R.L. Jaffe, Phys. Rev. **D9** (1974) 1444; (E) **D10** (1974) 1669.

[31] V. Bernard, N. Kaiser and U-G. Meissner, Eur. Phys. J **A4** (1999) 259.